

# More kaonic bound states and a comprehensive interpretation of the $D_{sJ}$ states

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## Abstract

The leading order interaction between a Goldstone boson and a matter field is universally dominated by the Weinberg-Tomozawa term. Based on this observation, we predict a rich spectrum of bound states of a kaon and a heavy meson. We argue that if the life time of an excited heavy meson is significantly longer than the range of forces, then the finite width of that state can be neglected in a first approximation. Then, the  $D_{s0}^*(2317)$ ,  $D_{s1}(2460)$ ,  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  are generated as  $DK$ ,  $D^*K$ ,  $D_1(2420)K$  and  $D^*(2600)K$  bound states, respectively. In addition to the remarkable agreement with the measured masses, the decay patterns of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  can also be understood. Two more  $D_{sJ}$  states, and kaonic bound states with the bottom mesons as well as the doubly charmed baryon are also predicted.

PACS numbers: 14.40.Lb, 12.39.Fe, 13.25.Ft

The leading order (LO) interaction between a Goldstone boson and a matter field is universally dominated by the Weinberg-Tomozawa (WT) term [1, 2]. With the kaon being a Goldstone boson of the spontaneous chiral symmetry breaking, the  $D_{s0}^*(2317)$  and  $D_{s1}(2460)$  can be generated as bound states of the  $DK$  and  $D^*K$  systems, respectively [3–9]. Considering a heavy matter field  $H$  with a mass  $M_H$  much larger than the kaon mass  $M_K$ , the heavy field can be treated nonrelativistically. Then the LO  $S$ -wave elastic scattering amplitude from the WT term can be given in a simple form,

$$V(s) = C \frac{M_H E_l(s)}{F_\pi^2}, \quad (1)$$

where  $s$  is the squared energy in the center-of-mass frame,  $E_l$  is the energy of the light meson in the center-of-mass frame, and  $F_\pi = 92.4$  MeV is the pion decay constant. The coefficient  $C = -2$  for the  $I = 0$   $HK \rightarrow HK$  channel. The heavy matter fields for which Eq. (1) can be applied include the excited heavy mesons with flavor contents  $c\bar{q}, b\bar{q}$ , and the doubly heavy antibaryons  $\bar{c}\bar{c}\bar{q}$  and  $\bar{b}\bar{b}\bar{q}$ . (A subtlety related to the finite width effects will be discussed later.) Therefore, if the  $D_{s0}^*(2317)$  is a  $DK$  bound state as suggested in Refs. [3, 4, 9, 10], one can easily expect that there must be more bound states. In fact, the  $\bar{B}K$  and  $\bar{B}^*K$  bound states have been predicted [3–5, 8]. Here a bound state does not strictly mean a state with a zero width. It means if all the other particles are neglected, then it has a vanishing width. If some other particles with lower masses are switched on as in the real world, these “bound states” will have a finite width.

On the experimental side, besides the  $D_{s0}^*(2317)$  and  $D_{s1}(2460)$ , more charmed strange mesons have been discovered in recent years. They include the  $D_{s1}^*(2700)$  [11, 12],  $D_{sJ}(2860)$  [11] and  $D_{sJ}(3040)$  [13]. Both the  $D_{s1}^*(2700)$  and the  $D_{sJ}(2860)$  decay into  $DK$  and  $D^*K$ , while the  $D_{sJ}(3040)$  was only observed in the  $D^*K$  final state [13]. The ratios between different decay modes of

the  $D_{sJ}(2860)$  were measured to be [14]

$$R_{D_{sJ}(2860)} = \frac{\Gamma(D_{sJ}(2860) \rightarrow D^*K)}{\Gamma(D_{sJ}(2860) \rightarrow DK)} = 1.10 \pm 0.24, \quad (2)$$

and a similar ratio for the  $D_{s1}^*(2700)$  is  $R_{D_{s1}^*(2700)} = 0.91 \pm 0.18$ . Both states have a natural parity, i.e. a positive (negative) parity for an even (odd) spin, because they decay into the  $DK$  channel. Using the LO heavy hadron chiral perturbation theory (HHChPT), which is model-independent, the decay pattern of the  $D_{s1}^*(2700)$  for a  $(2S, J^P = 1^-)$  assignment was calculated to be  $R_{D_{s1}^*(2700)} = 0.91 \pm 0.04$  [15], fully consistent with the data. Furthermore, the mass of the  $D_{s1}^*(2700)$  agrees well with the prediction in the Godfrey-Isgur quark model [16]. So there is little doubt that the  $D_{s1}^*(2700)$  is the  $(2S, 1^-)$   $c\bar{s}$  meson. The situation for the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  is not clear yet. Their quantum numbers are not known. The LO results from the HHChPT for  $R_{D_{sJ}(2860)}$  were given in [17] for different assignments. They are 1.23, 0, 0.63, 0.06, and 0.39 for  $(2S, 1^-)$ ,  $(2P, 0^+)$ ,  $(2P, 2^+)$ ,  $(1D, 1^-)$ , and  $(1D, 3^-)$ , respectively. Comparing with the measured value, one can conclude that the only possibilities are  $(2S, 1^-)$  and  $(2P, 2^+)$ . However, the  $(2S, 1^-)$   $c\bar{s}$  meson has already been identified as the  $D_{s1}^*(2700)$ , and the  $(2P, 2^+)$  one would have a mass as large as 3.1 GeV [18]. Various explanations of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  have already been discussed in the literature [19–27].

In this paper, we will show that all of the  $D_{s0}^*(2317)$ ,  $D_{s1}(2460)$ ,  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  mesons can be interpreted as charmed-meson kaon bound states. The  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  are interpreted as  $D_1(2420)K$  and  $D^*(2600)K$  bound states, respectively. In addition, more kaonic bound states will be predicted.

One important issue requires discussion. Both the  $D$  and  $D^*$  have a negligible width, however, the excited heavy mesons normally have a width of tens of MeV or even larger than 100 MeV. The finite width presents

another energy scale which might invalidate the use of the WT term. A natural question is: how large a width can be neglected for a given case? Let us consider a hadron with a finite life time  $\tau$ . If  $\tau$  is long enough so that the interaction responsible for the binding takes place, then one can neglect the width in calculating the bound state masses. Mathematically, this means  $\tau \gg r$ , with  $r$  the range of forces, i.e.  $\Gamma \ll 1/r$ . For the interaction between a kaon and a heavy meson, since a pion cannot be exchanged, the range of forces is set by  $1/r \sim M_\rho$  or more safely by  $2M_\pi$ . Therefore, if the width of the other hadron is much smaller than  $M_\rho$ , then in a first approximation, the width can be neglected. These heavy mesons, written in terms of  $\{H, H_s\}$  in the same SU(3) triplet, include  $\{D, D_s\}$ ,  $\{D^*, D_s^*\}$ ,  $\{D_1(2420), D_{s1}(2536)\}$ ,  $\{D_2^*(2460), D_{s2}^*(2573)\}$ ,  $\{D(2550), ?\}$ ,  $\{D^*(2600), D_{s1}^*(2700)\}$ , and  $\{\bar{B}, \bar{B}_s\}$ ,  $\{\bar{B}^*, \bar{B}_s^*\}$ ,  $\{B_1(5720), B_{s1}(5830)\}$ ,  $\{B_2^*(5747), B_{s2}(5840)\}$  [14], and their strange partners. Among them, the  $D(2550)$  and  $D^*(2600)$  were discovered very recently by the BABAR Collaboration [28]. The properties of the  $D^*(2600)$  are consistent with the radially excited  $(2S, 1^-)$  charmed meson. The  $D(2550)$  might be the  $(2S, 0^-)$  state since it has the correct angular distribution [28]. However, its width is much larger than the prediction using the LO HHChPT,  $\Gamma_{D(2550)} = 0.55\Gamma_{D^*(2600)} \sim 50$  MeV [29]. Nevertheless, we will tentatively assume the  $D(2550)$  as the  $2S$  pseudoscalar charmed meson. On the contrary, however, this is not the case for the  $\bar{D}_1(2430)D^*$  system which was considered in Ref. [30]. The width of the  $D_1(2430)$  is  $\sim 400$  MeV  $\gg M_\pi$ , whose inverse sets the range of forces for this system, and hence the finite width effect cannot be neglected. Indeed, it was shown that, after considering the finite width or the three-body cut, the bound state disappears [31]. The impact of the finite width effect on the line shape of a composite particle was considered in Ref. [32] (see also [33]).

The masses of the bound states can be calculated by searching for poles in the resummed  $S$ -wave  $I = 0$  scattering amplitudes  $T(s) = V(s)[1 - G(s)V(s)]^{-1}$  [34, 35]. We will consider two coupled channels for each case, which means all of  $T(s)$ ,  $V(s)$  and  $G(s)$  are  $2 \times 2$  matrices. Denoting the nonstrange (strange) heavy meson by  $H_{(s)}$ , the matrix elements of the symmetric matrix  $V(s)$  is given by Eq. (1). For  $I = 0$ , the coefficient  $C = -2, 0$  and  $-\sqrt{3}$  for  $V_{HK \rightarrow HK}(s)$ ,  $V_{H_s \eta \rightarrow H_s \eta}(s)$  and  $V_{HK \rightarrow H_s \eta}(s)$ , respectively.  $G(s)$  is a diagonal matrix with the nonvanishing elements given by the loop function for a nonrelativistic heavy meson and a relativistic light meson [8]

$$G_{Hl}(s) = \frac{1}{16\pi^2 M_H} \left\{ E_l \left[ a(\mu) + \log \left( \frac{M_l^2}{\mu^2} \right) \right] + 2|\vec{p}_l| \cosh^{-1} \left( \frac{E_l}{M_l} \right) - 2\pi i |\vec{p}_l| \right\}, \quad (3)$$

with  $a(\mu)$  a subtraction constant,  $\mu$  the scale of dimensional regularization, and  $\vec{p}_l(M_l)$  the three-momentum (mass) of the light meson. Equation (1) seems to imply that the heavier the matter field is, the stronger the interaction will be. However, the factor  $M_H$  will be canceled by a factor of  $1/M_H$  in the loop function in the resummed amplitudes, as required by the heavy quark symmetry. It is natural to choose the same value of the subtraction constant  $a(\mu)$  for all the heavy hadrons. A change in the scale  $\mu$  can be balanced by a corresponding change in  $a(\mu)$ . However, for checking the stability of the results, we will also fix the subtraction constant while allowing  $\mu$  varying from 1 GeV to  $M_H$ . For a given  $\mu$ ,  $a(\mu)$  is fixed by reproducing the mass of the  $D_{s0}^*(2317)$ , and we get  $a(1 \text{ GeV}) = -3.84$ . The results for the charmed mesons are summarized in Table I with the expected dominant decay modes. We use the central values for the masses of the constituents. The error bars reflect the variation of  $\mu$ . It has been assumed that the unknown strange partner of the  $D(2550)$  has a mass 100 MeV heavier than the  $D(2550)$ . In fact, neglecting the  $H_s \eta$  channel and keeping only the  $HK$  channel, the results are almost the same, with a difference within 4 MeV. This means one can interpret the generated states as heavy-meson kaon bound states. We can see that with input only from the  $D_{s0}^*(2317)$ , the masses of the  $D_{s1}(2460)$ ,  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  can be well reproduced. In the following, we will focus on the latter two states, which are interpreted as  $S$ -wave  $D_1(2420)K$  and  $D^*(2600)K$  bound states, respectively. In this scenario, the  $J^P$  are  $1^-$  for the  $D_{sJ}(2860)$ , and  $1^+$  for the  $D_{sJ}(3040)$ . We also checked that, using  $F_K$  instead of  $F_\pi$  in Eq. (1), which represents some of the higher order corrections in the chiral expansion, the results are very stable. The predicted masses change for no more than 3 MeV with the subtraction constant refitted to the mass of the  $D_{s0}^*(2317)$ . Inclusion of other coupled channels such as those listed as decay modes in Table I involves unknown coupling constants. As long as the threshold of the coupled channel  $T_{CC}$  is far from the bound state mass, which can be characterized as  $|T_{CC} - M_{BS}| \gg \epsilon$ , with  $\epsilon$  being the binding energy, they will modify the mass only marginally. However, their presence will give a finite width to the bound state. Now let us discuss their decay modes.

Some of the decay mechanisms of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  are shown in Fig. 1. Note that this is not a complete list, for instance, the  $D_s \omega$  is another important decay channel of the  $D_{sJ}(3040)$ . In fact, this channel can be used in distinguishing the  $D^*(2600)K$  bound state picture from a  $c\bar{s}$  picture, for which the decay into  $D_s \omega$  would be very small due to the OZI suppression. Being a  $1^+$  state, the  $D_{sJ}(3040)$  cannot decay into the  $DK$ . This is consistent with the experimental facts. The two-body decays of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  occur through the scattering  $D_1(2420)K \rightarrow D^{(*)}K$  and

TABLE I: Predicted masses of charmed-meson kaon bound states. The experimental values are listed in the fourth row for a comparison. The expected dominant decay modes are also given.

Main constituents	$DK$	$D^*K$	$D_1(2420)K$	$D_2(2460)K$	$D(2550)K$	$D^*(2600)K$
$J^P$	$0^+$	$1^+$	$1^-$	$2^-$	$0^+$	$1^+$
Predicted masses	2317.8(input)	$2458 \pm 3$	$2870 \pm 9$	$2910 \pm 9$	$2984 \pm 10$	$3052 \pm 11$
Experimental data	$2317.8 \pm 0.6$	$2459.5 \pm 0.6$	$2862 \pm 2_{-2}^{+5}$			$3044 \pm 8_{-5}^{+30}$
Dominant decays	$D_s\pi$	$D_s^*\pi$	$D^{(*)}K, D_s^{(*)}\eta$	$D^*K, D_s^*\eta$	$DK, D_s\eta$	$D^*K, D_s^*\eta, D_s\omega, DK^*, D\phi$

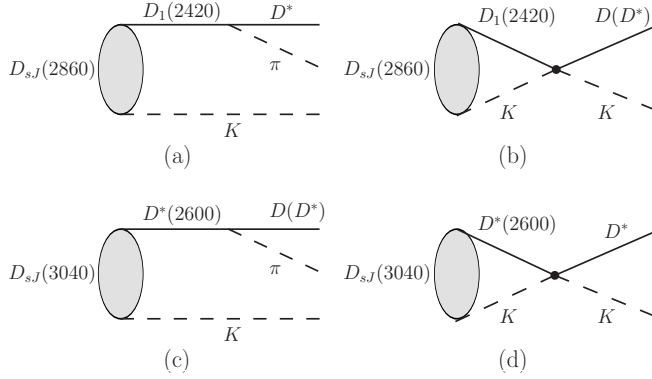


FIG. 1: Schematic diagrams of some of the decay mechanisms (not a complete list) of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$ .

$D^*(2600)K \rightarrow D^*K$ , respectively, as shown in Fig. 1. The total angular momentum of the light degrees of freedom in a heavy hadron  $s_\ell$  is conserved in the heavy quark limit. The axial-vector meson  $D_1(2420)$  has  $s_\ell^P = \frac{3}{2}^+$ . So without breaking the heavy quark spin symmetry, the  $D_1(2420)K$  can couple to the  $DK$  and  $D^*K$  channels in a  $P$  wave because  $D^{(*)}$  has  $s_\ell^P = \frac{1}{2}^-$ . Therefore, the two-body decay modes  $D^*K$  and  $DK$  can be related through heavy quark spin symmetry. Nonrelativistically, the decay amplitudes can be written as

$$\begin{aligned} \mathcal{M}(D_{sJ}(2860) \rightarrow DK) &= g_D G_{D_1K} \vec{\epsilon}_{D_{sJ}} \cdot \vec{k}_D, \\ \mathcal{M}(D_{sJ}(2860) \rightarrow D^*K) &= g_{D^*} G_{D_1K} \epsilon^{ijk} k_D^i k_{D^*}^j \vec{\epsilon}_{D_{sJ}}^k \vec{\epsilon}_{D^*}^k \end{aligned} \quad (4)$$

where  $G_{D_1K}$  is the  $D_1K$  loop function,  $\vec{k}_{D^{(*)}}$  is the momentum of the  $D^{(*)}$  in the center-of-mass frame,  $\epsilon^{ijk}$  is the antisymmetric Levi-Civita tensor, and  $\vec{\epsilon}_{D_{sJ}}$  and  $\vec{\epsilon}_{D^*}$  are the polarization vectors of the  $D_{sJ}(2860)$  and  $D^*$ , respectively. The heavy quark spin symmetry requires  $g_D/g_{D^*} = \sqrt{M_D/M_{D^*}}$  taking into account the nonrelativistic normalization factor. The result

$$R_{D_{sJ}(2860)} = 2 \frac{M_{D^*}}{M_D} \left| \frac{\vec{k}_{D^*}}{\vec{k}_D} \right|^3 = 1.23 \quad (5)$$

is in nice agreement with the data  $1.10 \pm 0.24$ . Furthermore, the helicity angular distribution for the sequential process  $D_{sJ}(2860) \rightarrow D^*K \rightarrow D\pi K$  derived from Eq. (4) is  $\sin^2 \theta$ , which again agrees with the observation [13].

One expects that the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  can also decay through the decays of the  $D_1(2420)$  and  $D^*(2600)$ , respectively. However, the partial widths of these sequential decays are not large. This is because they are suppressed by a  $D$ -wave and  $P$ -wave factor for the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$ , respectively, as well as by the three-body phase space. In fact, this expectation can be confirmed by a rough estimate. For an  $S$ -wave loosely bound state, the binding energy  $\epsilon$  and effective coupling  $g$  for the bound state to its constituents are related by  $g^2 = 16\pi(m_1 + m_2)^2 \sqrt{2\epsilon/\mu} [1 + \mathcal{O}(r\sqrt{2\mu\epsilon})]$  [36, 37], where  $m_{1,2}$  are the masses of the constituents, and  $\mu$  is their reduced mass. We get 1 and 26 MeV for such sequential decays of the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$ , respectively, which means they only contribute a branching fraction of about 2% and 10%, respectively. However, these results can only be regarded as an order-of-magnitude estimate, since for both cases,  $\sqrt{2\mu\epsilon} \simeq 220$  MeV, and so the uncertainty in the so-estimated effective coupling is very large. Nevertheless, the two-body decays would be the dominant modes for these two states. Therefore, one expects  $\Gamma_{D_{sJ}(3040)} \gg \Gamma_{D_{sJ}(2860)}$ , consistent with the data, because the two-body decays of the  $D_{sJ}(3040)$  and  $D_{sJ}(2860)$  are in an  $S$ - and  $P$ -wave, respectively.

In addition, heavy quark spin symmetry implies that each of the predicted bound states has its spin multiplet partner [38]. For instance, the  $D_{s1}(2460)$  is the spin partner of the  $D_{s0}^*(2317)$ . Because the kaon, which has a negative parity, interacts with the heavy meson in an  $S$ -wave, so for the  $D^{(*)}K$  systems,  $s_\ell^P = \frac{1}{2}^+$ . Note that they should not be confused with the  $s_\ell^P = \frac{1}{2}^+ c\bar{s}$  mesons because they have a different dynamical origin. Similarly, the  $D_2(2460)K$  bound state, to be called  $D_{s2}^*(2910)$ , is the spin partner of the  $D_{sJ}(2860)$ , and they have  $s_\ell^P = \frac{3}{2}^-$ . The  $D_{sJ}(3040)$  and the  $D(2550)K$  bound state, to be called  $D_{s0}^*(2985)$ , form another  $s_\ell^P = \frac{1}{2}^+$  doublet. They can be regarded as the excited states of the  $D_{s0}^*(2317)$  and  $D_{s1}(2460)$ .

Assuming the decay modes given in Table I exhaust the decay widths, ratios of the total widths can be predicted based on the HHChPT. At LO, we obtain

$$\frac{\Gamma_{D_{s2}^*(2910)}}{\Gamma_{D_{sJ}(2860)}} \simeq 0.2, \quad \frac{\Gamma_{D_{s0}^*(2985)}}{\Gamma_{D_{sJ}(3040)}} < 1, \quad (6)$$

where the less-than sign is because the  $DK^*, D_s\omega, D_s\phi$

channels for the  $D_{sJ}(3040)$  were not taken into account. Thus we have  $\Gamma_{D_{s2}^*(2910)} \sim 10$  MeV. Such a small width suggests that more experimental efforts should be devoted to the  $D^*K$  data.

We can make further predictions for the bottom sector. The results are listed in Table II together with the expected dominant decay channels. In addition, there

TABLE II: Predicted masses of bottom-meson kaon bound states. The expected dominant decay modes are also given.

Constituents	$\bar{B}K$	$\bar{B}^*K$	$\bar{B}_1(5720)K$	$\bar{B}_2(5747)K$
$J^P$	$0^+$	$1^+$	$1^-$	$2^-$
Predicted masses	$5705 \pm 31$	$5751 \pm 32$	$6151 \pm 33$	$6169 \pm 33$
Dominant decays	$\bar{B}_s\pi$	$\bar{B}_s^*\pi$	$\bar{B}^{(*)}K, \bar{B}_s^{(*)}\eta$	$\bar{B}^*K, \bar{B}_s^*\eta$

should also be kaonic bound states with the doubly heavy baryons. For instance, the  $\bar{\Xi}_{cc}(3520)K$  bound state is predicted to have a mass of  $3956 \pm 20$  MeV.

In summary, using chiral and heavy quark symmetry, we predicted a rich spectrum of kaonic bound states. We argue that, if the lifetime of the constituents of a bound state is significantly longer than the range of forces, then the finite width effect can be neglected in a first approximation. The  $D_1(2420)K$  and  $D^*(2600)K$  bound states are found to fit very well to both the measured masses and decay patterns for the  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$ , respectively. Hence, all known properties of the  $D_{s0}^*(2317)$ ,  $D_{s1}(2460)$ ,  $D_{sJ}(2860)$  and  $D_{sJ}(3040)$  can be systematically explained as kaonic bound states. Kaonic bound states for the bottom mesons and doubly charmed baryon are also predicted. The decay mode  $D_{sJ}(3040) \rightarrow D_s\omega$  can serve as a criterion in distinguishing the present interpretation from a  $c\bar{s}$  picture. The narrow  $D_{s2}^*(2910)$  and broader  $D_{s0}^*(2985)$  should be searched for in the  $D^*K$  and  $DK$  channels, respectively. In order to understand the low-energy strong interaction better, searching for these states should be an important experimental issue.

We are grateful to C. Hanhart for useful comments. We thank the HGF for funds provided to the virtual institute “Spin and strong QCD” (VH-VI-231), the DFG (SFB/TR 16) and the EU I3HP “Study of Strongly Interacting Matter” under the Seventh Framework Program of the EU. U.-G. M. also thanks the BMBF for support (Grant No. 06BN9006).

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